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Scattering of strong electromagnetic wave by relativistic electrons: Thomson and Compton regimes

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ABSTRACT

The processes of the nonlinear Compton and the nonlinear Thomson scattering in a field of intense plane electromagnetic wave in terms of photon yield have been considered. The quantum consideration of the Compton scattering process allows us to calculate the probability of a few successive collisions k of an electron with laser photons accompanied by the absorption of n photons (nonlinear regime) when the number of collisions and the number of absorbed photons are of random quantities. The photon spectrum of the nonlinear Thomson scattering process was obtained from the classical formula for intensity using the Planck's law. The conditions for which the difference between the classical and the quantum regimes is manifested was obtained. Such a condition is determined by a discrete quantum radiation mechanism, namely, by the mean number of photons $\bar{k}$ emitted by an electron passing through the laser pulse.

1. Introduction

During the last years, a few accelerator facilities have been developing new sources of mono-chromatic X- and γ-rays which are based on the Compton/Thomson backscattering (CBS/TS) processes [1–17]. One of the main purposes of the ELI-NP project is to produce an intense γ-beam with the energy $\hbar\omega_0 \sim 10$ MeV and a monochromaticity much less than 1% [18] using tight collimation with an opening angle $\theta_c \ll \gamma_0^{-1}$ (γ_0 is the Lorentz-factor of initial electrons). In order to get the intensity of such a collimated γ-beam which can also be used for different applications, there should be used an intense laser pulse or, in other words, “the strong wave” which can be characterized by a value of the dimensionless laser strength parameter a_0:

$$a_0 = \frac{e}{mc}\frac{E_0}{\omega_0}. \quad (1)$$

In (1), e and m are the charge and the mass of an electron, c is the speed of light, E_0 is the laser field, ω_0 is the radiation frequency. For the strong wave, such a parameter can achieve the value compared with unity (or more).

The well-known process of the intense laser radiation scattering by free electrons in the classical electrodynamics is described as the nonlinear Thomson scattering (TS). The same process in the quantum electrodynamics is treated as the nonlinear Compton scattering (CS). In the former case, the electron oscillates in the laser wave and continuously emits radiation consisting of n harmonics ($n \geq 1$). In the latter case there is a discrete process of the photon emission in which an electron “absorbs” $n \geq 1$ laser photons with an energy $\hbar\omega_0$ and emits “hard” photon only with an energy $\hbar\omega \sim n\hbar\omega_0$. Many authors have considered both radiation mechanisms (see, for instance, [19–22] and cited papers there).

The drastic difference between TS and CS regi- mes occurs for ultrashort strong laser wave $a_0 \geq 1$, $N_0 \sim 10^1$, N_0 is the number of cycles in the laser pulse (see, for instance, [19,20]). For “long” laser pulses ($a_0 \sim 1$, $N_0 \geq 10^2$) the spectral distributions of emitted photons for both mechanisms are similar and the difference is in the so-called “red-shift”. Such a red-shift characterizes the decrease of the energies of the spectral maximum for the nonlinear Compton scattering in comparison with the Thomson scattering due to the quantum recoil effect.

The authors of the paper [21] compared both nonlinear mechanisms and obtained the following conditions providing a coincidence of results from the Compton and the Thomson mechanisms:

$$n\frac{4\gamma_0\hbar\omega_0}{mc^2(1 + a_0^2)} \ll 1. \quad (2)$$

Here, n is the number of absorbed photon, $\hbar\omega_0$ is the energy of the laser photon.

It means that an energy of the laser photon in the rest system

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(where the electron is in a rest in average) is much less than the electron mass and, accordingly, a quantum recoil effect does not change the energy of scattered photons practically.

In the paper [21] the authors claimed that for ultrashort femtosecond laser pulses and $\gamma_0 \leq 10^3$, $a_0 \sim 1$ "the classical and quantum results are practically the same" if the condition $\gamma_0 \hbar\omega_0/mc^2 \ll 1$ is fulfilled.

Seipt and Kämpfer in the work [19] found that spectral distributions for radiation at a fixed angle consisting of a set of sub-peaks are determined by the temporal shape of the ultrashort laser pulse and differ for both mechanisms. The energy of scattered photons will be strictly defined by the emission angle and the parameter a_0 only for long pulses ($N_0 \gg 10^2$). They showed that Compton and Thomson cross-sections practically coincide if the condition (2) is fulfilled. Also, they noticed that for $a_0 \gg 1$ these cross-sections will be different although the above-mentioned condition is kept.

In our work, we concentrated our attention on the discrete process of photon emission which can't be considered in the classical approach. Below we shall show that the process of emission of a few photons by each electron during successive interactions with laser photons (multi-photon emission [24] or multiple Compton backscattering [25]) can only be considered in the quantum approach. Such a multiple back-scattering process leads to a significant difference between the Compton and the Thomson spectra even if the condition Eq. (2) is fulfilled.

We consider the multiple Compton backscattering process as the independent emission of photons because the formation length of the emitted photon $\sim\gamma^2\lambda$ [26] is much less than a length of the laser pulse $c\tau$. For an electron energy $\sim 10^3$ MeV and an energy of the emitted photon ~ 10 MeV ($\lambda \sim 10^{-7}$ um) the formation length achieves macroscopic size 0.5 μm but for a laser pulse duration $\tau \sim 1$ ps the pulse length is much longer than the formation length.

The paper is organized as follows: the Section 1 is devoted to the Introduction. In Section 2, we consider the nonlinear Thomson scattering process in the classical electrodynamics frame and obtain the semi-classical emitted photon spectrum from the classical expression for the intensity spectrum using the Planck's law. The consideration of the nonlinear Compton backscattering process as the typical quantum process was performed in Section 3 on the base of formulas for total cross-sections and for the luminosity. The differential cross-sections characterizing the nonlinear Compton scattering were obtained in Section 4 for small values of the laser strength parameter a_0 and by taking into account the number of absorbed laser photons $n = 1, 2, 3, \ldots$. In Section 5 we describe the Monte-Carlo algorithm for simulation of the nonlinear Compton scattering process where an electron absorbs n laser photons and in each collision emits $k = 0, 1, 2, \ldots$ hard quanta (multiple Compton backscattering). The main results obtained after comparing both approaches considered in the paper, are discussed in Section 6. The final conclusions are presented in Section 7.

2. Photon flux from the nonlinear Thomson scattering process

Let us write the known formulas for the nonlinear TS for a strong circularly-polarized laser wave scattering on a counterpropagated relativistic electron. The incoherent sum of two circularly polarized waves with equal intensities and opposite helicities gives the result, coinciding with an unpolarized beam. So our consideration can be extended for unpolarized radiation because a spectral-angular distribution of scattered photons does not depend on a helicity of the initial wave.

Below, we consider the radiation of an electron in the field of the strong monochromatic wave only. In this approach, one can use the well-known relation connecting the frequency of n-th harmonic and emission angle [20,22].

$$\omega^{(n)} = n\frac{4\gamma_0^2\omega_0}{1 + (\gamma_0\theta)^2 + a_0^2/2}. \tag{3}$$

It is convenient to use a dimensionless spectral variable:

$$S^{(n)} = \frac{\omega^{(n)}}{4\gamma_0^2\omega_0} = \frac{n}{1 + (\gamma_0\theta)^2 + a_0^2/2}, \quad 0 \leq S^{(n)} \leq S_{\max}^{(n)},$$
$$S_{\max}^{(n)} = n/(1 + a_0^2/2). \tag{4}$$

The spectral distribution of emitted photons of an electron passing the laser field for the fixed harmonic number n can be obtained from the classical intensity spectrum dividing one of the emitted photon energy $\hbar\omega$ [22].

$$\frac{dN^{(n)}}{dS^{(n)}} = 2\pi\alpha a_0^2 N_0 \times \left\{ \frac{[n - S^{(n)}(2 + a_0^2)]^2}{2S^{(n)}a_0^2[n - S^{(n)}(1 + a_0^2/2)]} J_n^2(nz) + J_{n'}^2(nz) \right\}. \tag{5}$$

In the formula (5) N_0 is the number of electron oscillations along its trajectory governed by a laser field, the Bessel function and its derivative of order n, are denoted by $J_n(x)$, $J'_n(x)$, were argument nz is

$$nz = \sqrt{2}a_0\sqrt{S^{(n)}n - S^{(n)^2}(1 + a_0^2/2)}\,.$$

The number of photons emitted at the n-th harmonic can be calculated after integration of the spectrum (5):

$$N^{(n)} = \int_0^{S_{\max}^{(n)}} \frac{dN^{(n)}}{dS^{(n)}} dS^{(n)}. \tag{6}$$

It is evidently that such a value is determined by the field strength parameter a_0 and the number of cycles only. In Fig. 1 we show the dependence of the emitted photons number $N^{(n)}$ per an electron and per an oscillation ($N_0 = 1$) on the parameter a_0 (for n=1, 2, 3). The total number of emitted photons can be found by summing up the contributions from the n-th order harmonic number (6).

$$N_{\text{tot}} = \sum_{n=1}^{\infty} N^{(n)}. \tag{7}$$

In order to realize the summation procedure, the maximal harmonic number $n_{\max}$, which depends on the parameter a_0, should be chosen. For this purpose we suggest the following criterion:

$$\frac{N^{(n_{\max}+1)}}{N_{\text{tot}}} < 0.0005, \tag{8}$$

where

$$N_{\text{tot}} = \sum_{n=1}^{n_{\max}} N^{(n)}. \tag{9}$$

Such a choice provides a relative accuracy in the calculation of the value N_{tot} not more than 1–2%. The dependence of the total number of emitted photons N_{tot} on the laser field strength is shown in Fig. 2, where

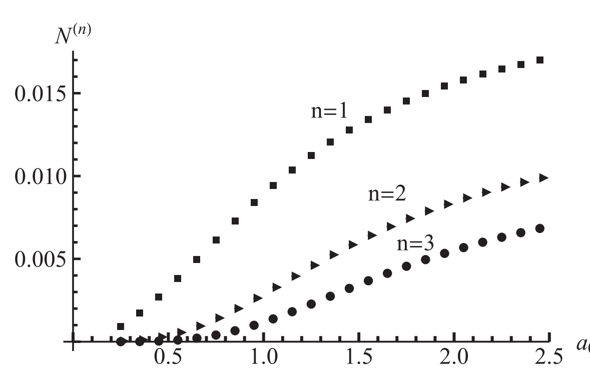

Fig. 1. Dependences of the emitted photon number $N^{(n)}$ for the different harmonic number ($n = 1, 2, 3$) on the strength field parameter a_0 per a cycle of the wave.